

Systematic Study of High p_T Hadron Spectra in pp , pA and AA Collisions from SPS to RHIC Energies

Xin-Nian Wang

Nuclear Science Division, Mailstop 70A-3307,
Lawrence Berkeley National Laboratory, Berkeley, CA 94720 USA
(November 17, 1998)

High- p_T particle spectra in $p + p$ ($\bar{p} + p$), $p + A$ and $A + B$ collisions are calculated within a QCD parton model in which intrinsic transverse momentum, its broadening due to initial multiple parton scattering, and jet quenching due to parton energy loss inside a dense medium are included phenomenologically. The intrinsic k_T and its broadening in $p + A$ and $A + B$ collisions due to initial multiple parton scattering are found to be very important at low energies ($\sqrt{s} < 50$ GeV). Comparisons with $S + S$, $S + Au$ and $Pb + Pb$ data with different centrality cuts show that the differential cross sections of large transverse momentum pion production ($p_T > 1$ GeV/c) in $A + B$ collisions scale very well with the number of binary nucleon-nucleon collisions (modulo effects of multiple initial scattering). This indicates that semi-hard parton scattering is the dominant particle production mechanism underlying the hadron spectra at moderate $p_T \gtrsim 1$ GeV/c. However, there is no evidence of jet quenching or parton energy loss. Within the parton model, one can exclude an effective parton energy loss $dE_q/dx > 0.01$ GeV/fm and a mean free path $\lambda_q < 7$ fm from the experimental data of $A + B$ collisions at the SPS energies. Predictions for high p_T particle spectra in $p + A$ and $A + A$ collisions with and without jet quenching at the RHIC energy are also given. Uncertainties due to initial multiple scattering and nuclear shadowing of parton distributions are also discussed.

PACS numbers: 25.75.-q, 12.38.Mh, 13.87.-a, 24.85.+p

I. INTRODUCTION

Large- E_T partons or jets are good probes of the dense matter formed in ultra-relativistic heavy-ion collisions [1,2], since they are produced in the earliest stage of heavy-ion collisions and their production rates are calculable in perturbative QCD. If a dense partonic matter is formed during the initial stage of a heavy-ion collision with a large volume and a long life time (relative to the confinement scale $1/\Lambda_{\text{QCD}}$), the produced large E_T parton will interact with this dense medium and, according to many recent theoretical studies [3-5], will lose its energy via induced radiation. The energy loss is shown to depend on the parton density of the medium. Therefore, the study of parton energy loss can shed light on the properties of the dense matter in the early stage of heavy-ion collisions. It is also a crucial test whether there is any thermalization going on in the initial stage of heavy-ion collisions.

Even though one cannot measure directly the energy loss suffered by a high energy parton propagating through a dense medium, its effective fragmentation functions must change and leading particles must be suppressed due to the parton energy loss. This can be measured directly in deeply inelastic $e + A$ collisions (for jets going through a normal nuclear matter) [6,7] and direct-photon-tagged jets in high-energy heavy-ion collisions [7,8]. Since the large p_T single-inclusive particle spectra in nuclear collisions are a direct consequence of jet

fragmentation as in hadron-hadron and hadron-nucleus collisions, they are also shown [9] to be sensitive to parton energy loss or modification of jet fragmentation functions inside a dense medium.

Because of the extremely small cross sections of large p_T particle production at the CERN SPS energy ($\sqrt{s} \approx 20$ GeV), measurements of particle spectra in heavy-ion collisions at moderately large p_T have not become available before recent experiments [10,11]. In a recent paper [12], the author used a QCD parton model calculations to analyze the large p_T spectra and found that there is no evidence of parton energy loss in high-energy heavy-ion collisions. From the analysis of the spectra, one can limit the effective energy loss to less than $dE/dx < 0.01$ GeV/fm. Since experimental data have shown evidences of an interacting hadronic matter in heavy-ion collisions at the CERN SPS energy, one can also conclude that the hadronic matter does not cause apparent jet energy loss.

In this paper, we will conduct a systematic study of high p_T hadron spectra in pp , pA and AB collisions, since the reliability of the determination of a small effective parton energy loss in high-energy heavy-ion collisions crucially depends on the precision to which we understand the spectra without parton energy loss. We will include both the intrinsic k_T in $p + p$ collisions and k_T broadening due to multiple initial-state scattering in $p + A$ and $A + A$ collisions in a phenomenological manner, which we find necessary to describe large p_T particle

production at energies below $\sqrt{s} < 50$ GeV. We will then compare the calculations with the experimental data for $S + S$, $S + Au$ and $Pb + Pb$ collisions at the CERN SPS with different centrality cuts and verify the scaling behavior characteristic of hard processes. Using the same parton model we will then predict the high p_T spectra in pA collisions at the RHIC energy and then calculate same spectra in AA collisions with and without effects of parton energy loss. The goal of this paper is to provide a practical and phenomenological description of the high p_T hadron spectra in $p + p$ and $p + A$ collisions. This will provide a baseline prediction for high-energy $A + A$ collisions against which we can study the effects of the parton energy loss caused by the dense medium.

II. $P + P$ COLLISIONS: INITIAL K_T

Large p_T particle production in high-energy hadron-hadron collisions has been shown to be a good test of the QCD parton model [13]. Assuming partons inside the colliding hadrons only have longitudinal momenta, the large p_T particle production cross section can be calculated as a convolution of elementary parton-parton scattering cross sections, parton distributions inside the hadron and parton fragmentation functions. The factorization theorem ensures that the parton distributions and fragmentation functions are universal and can be measured in other hard processes, *e.g.*, e^+e^- annihilation, deeply inelastic ep collisions. With high precision data, parton distributions and fragmentation functions have been parameterized including their QCD evolution with the momentum scale Q of the hard processes [14,15] which in turn can be used to calculate particle production in many other hard processes. This parton model with collinear approximation has been rather successful so far in describing large p_T particle production in high-energy $p\bar{p}$ collisions [16] with $\sqrt{s} > 50$ GeV.

The above described collinear parton model fails to account the experimental data on angular correlation of produced heavy quarks and the total transverse momentum distribution of the heavy quark pairs [18] or the Drell-Yan lepton pairs [17]. In terms of parton models one naturally expects that partons inside a hadron carry at least an average intrinsic transverse momentum of about a few hundred MeV, reflecting the hadron size via the uncertainty principle. Furthermore, higher order pQCD processes like initial state radiation or $2 \rightarrow 3$ subprocesses with additional radiated gluons can also give rise to large initial transverse momentum for the colliding partons. However, even within next leading order (NLO) pQCD collinear parton model, experimental data of Drell-Yan lepton pairs production [17], heavy quark pair production [19], and direct photon production [20] all indicate that an average initial parton transverse momentum of about 1 GeV is needed to describe the data. Such a large value is an indication of its perturbative

nature. Theoretical techniques have been developed to include all high order contributions [21–23], though it is still problematic to differentiate what is true intrinsic and what is pQCD generated initial transverse momentum. One can re-sum all the radiative corrections into a “Sudakov” form factor up to a scale Q_0 and define (non-perturbative) contributions below this scale as “true” intrinsic, even though such an “intrinsic” average transverse momentum still depend on colliding energy \sqrt{s} and the momentum scale Q of the hard processes. Such schemes have been developed for processes, such as Drell-Yan [21–23] and heavy quark production [24] where two scales are involved, with only minimum phenomenological success in limited kinematic region. The problem in processes such as direct photon and single inclusive hadron production, where there is only a single scale, is even more difficult [25].

In this paper, we will adopt a more phenomenological approach by introducing initial parton transverse momentum with a Gaussian form of distribution as has been done in Drell-Yan [26], heavy quark [19], direct photon [20] and single inclusive hadron production [20,27,28]. Since it is difficult to extend the NLO calculation to $p + A$ and $A + A$ collisions where multiple collisions occur and our main purpose is to provide an effective description of single inclusive hadron spectra not only in $p + p$ but also in $p + A$ and $A + A$ collisions, we will use lowest order (LO) parton model with a simple scheme of including the intrinsic transverse momentum with a Gaussian distribution. The inclusion of initial transverse momentum can be considered as a parameterization of both higher order and nonperturbative corrections. The predictive power of this model then lies in the energy and flavor dependence of the hadron spectra. Similar comments can be made for our modeling of multiple parton scattering effect in $p + A$ and $A + A$ collisions.

One can define parton distributions in both fractional longitudinal momentum x and initial transverse momentum k_T in a factorized form,

$$dx d^2 k_T g_N(k_T, Q^2) f_{a/N}(x, Q^2), \quad (1)$$

where $g_N(k_T, Q^2)$ is the transverse momentum distribution and $f_{a/N}(x, Q^2)$ are the normal parton distribution functions for which we will use the MRS D-’ parameterization [14]. The inclusive particle production cross section in pp collisions will then be given by [13]

$$\begin{aligned} \frac{d\sigma_{pp}^h}{dy d^2 p_T} = & K \sum_{abcd} \int dx_a dx_b d^2 k_{aT} d^2 k_{bT} g_p(k_{aT}, Q^2) \\ & \times g_p(k_{bT}, Q^2) f_{a/p}(x_a, Q^2) f_{b/p}(x_b, Q^2) \\ & \times \frac{D_{h/c}^0(z_c, Q^2)}{\pi z_c} \frac{d\sigma}{dt}(ab \rightarrow cd), \end{aligned} \quad (2)$$

where $D_{h/c}^0(z_c, Q^2)$ is the fragmentation function of parton c into hadron h as parameterized in [15,29] from e^+e^- data, z_c is the momentum fraction of a parton jet carried by a produced hadron. The $K \approx 2$ (unless otherwise

specified) factor is used to account for higher order QCD corrections to the jet production cross section [30].

We define the momentum fraction x in terms of the light-cone variables $x_a = (E_a + p_{\parallel a})/\sqrt{s}$ for partons in the forward beam direction and $x_b = (E_b - p_{\parallel b})/\sqrt{s}$ in the backward beam direction. The four-vector momenta for the colliding partons are then $p_a = (E_a, \mathbf{k}_{Ta}, p_{\parallel a})$, $p_b = (E_b, \mathbf{k}_{Tb}, p_{\parallel b})$ with

$$\begin{aligned} E_a &= (x_a \sqrt{s} + \frac{k_{Ta}^2}{x_a \sqrt{s}})/2, \\ p_{\parallel a} &= (x_a \sqrt{s} - \frac{k_{Ta}^2}{x_a \sqrt{s}})/2, \\ E_b &= (x_b \sqrt{s} + \frac{k_{Tb}^2}{x_b \sqrt{s}})/2, \\ p_{\parallel b} &= (-x_b \sqrt{s} + \frac{k_{Tb}^2}{x_b \sqrt{s}})/2. \end{aligned} \quad (3)$$

In principle, one should also include the transverse momentum smearing from the jet fragmentation. We neglect this in our calculation of particle spectra in the central rapidity region and consider it been effectively included in the calculation by adjusting the initial k_T distribution. In this case, particles are produced in the same direction of the fragmenting jet. The Mandelstam variables for the elementary parton-parton scattering processes are then,

$$\begin{aligned} \hat{s} &= x_a x_b \sqrt{s} + \frac{k_{Ta}^2 k_{Tb}^2}{x_a x_b \sqrt{s}} - 2k_{Ta} k_{Tb} \cos(\phi_a - \phi_b), \\ \hat{t} &= -\frac{p_T}{z_c} \left(x_a \sqrt{s} e^{-y} + \frac{k_{Ta}^2}{x_a \sqrt{s}} e^y - 2k_{Ta} \cos \phi_a \right), \\ \hat{u} &= -\frac{p_T}{z_c} \left(x_b \sqrt{s} e^y + \frac{k_{Tb}^2}{x_b \sqrt{s}} e^{-y} - 2k_{Tb} \cos \phi_b \right), \end{aligned} \quad (4)$$

where p_T and y are the transverse momentum and rapidity of the produced particle, $\cos \phi_a = \mathbf{k}_{Ta} \cdot \mathbf{p}_T / k_{Ta} p_T$, and $\cos \phi_b = \mathbf{k}_{Tb} \cdot \mathbf{p}_T / k_{Tb} p_T$. The momentum fraction of the produced hadron z_c is then given by the identity for massless two-body scattering, $\hat{s} + \hat{t} + \hat{u} = 0$.

In Eq. (2), one should also restrict the initial transverse momentum $k_{Ta} < x_a \sqrt{s}$ and $k_{Tb} < x_b \sqrt{s}$ such that the partons' longitudinal momenta in Eq. (3) have the same signs as their parent hadrons. In addition, one of the Mandelstam variables can approach to zero if the initial k_T is too large and then the parton cross section could diverge. To avoid these problems, we introduce a regulator μ^2 in the denominators of the parton-parton cross sections $d\sigma/d\hat{t}(ab \rightarrow cd)$. We choose $\mu = 0.8$ GeV in our following calculations. The resultant spectra are sensitive to the choice of μ only at p_T around μ where we believe that QCD parton model calculation becomes unreliable.

We assume the initial k_T distribution $g_N(k_T)$ to have a Gaussian form,

$$g_N(k_T, Q^2) = \frac{1}{\pi \langle k_T^2 \rangle_N} e^{-k_T^2 / \langle k_T^2 \rangle_N}. \quad (5)$$

Since our initial k_T includes both the intrinsic and pQCD radiation-generated transverse momentum, the variance in the Gaussian distribution $\langle k_T^2 \rangle_N$ should depend on the momentum scale Q of the hard processes. As pointed out by Owen and Kimel [28], in addition to the leading log contribution from initial-state radiation that is included in the Q^2 dependence of the parton distributions, the recoil effect of large-angle gluon emissions can lead to Q -dependent $\langle k_T^2 \rangle_N$ in this approach of QCD parton model. In the following we shall consider both a constant $\langle k_T^2 \rangle_N$ and one that depends on the momentum scale of the hard processes. For the Q -dependent case, we choose the following form [28] in our model

$$\langle k_T^2 \rangle_N(Q^2) = 1.2(\text{GeV}^2) + 0.2\alpha_s(Q^2)Q^2. \quad (6)$$

The form of the Q -dependence and the parameters are chosen to reproduce the experimental data. Following the same approach as in Refs. [27,28], we choose Q^2 to be $Q^2 = 2\hat{s}\hat{t}\hat{u}/(\hat{s}^2 + \hat{t}^2 + \hat{u}^2)$.

Shown in Figs. 1-4 are our calculated spectra for charged pions as compared to the experimental data [31,32] for $p + p$ collisions at $E_{\text{lab}} = 200, 300, 400, 800$ GeV. As one can see from the figures that QCD parton model calculations with the initial k_T smearing (solid lines) as given in Eq. 6 fit the experimental data very well over all energy range. However, without the initial k_T smearing (dot-dashed lines for π^-) the calculations significantly underestimate the experimental data, as much as a factor of 20 at $E_{\text{lab}} = 200$ GeV. This is because the QCD spectra are very steep at low energies and even a small amount of initial k_T could make a big increase to the final spectra. As the energy increases, the QCD spectra become flatter and small amount of initial k_T does not change the spectra much, as we can already see by comparing the spectra for $E_{\text{lab}} = 800$ GeV (Fig. 4) to that of 200 GeV (Fig. 1). This is further demonstrated in Fig. 5 where we compare the LO calculation to experimental data in high-energy $p\bar{p}$ collisions. Here we have used a $K = 1.5$ factor. One can see that at higher collider energies, the initial k_T does not make much difference to the spectra at high transverse momentum. In the case of no intrinsic k_T , we didn't use the regulator μ in our calculation. Instead, we used the p_T of particle as a cut-off of the phase integral in Eq. (2). This is why the spectra (dot-dashed lines) increase faster at smaller p_T than the ones with initial k_T smearing in which a regulator of $\mu = 0.8$ GeV is used.

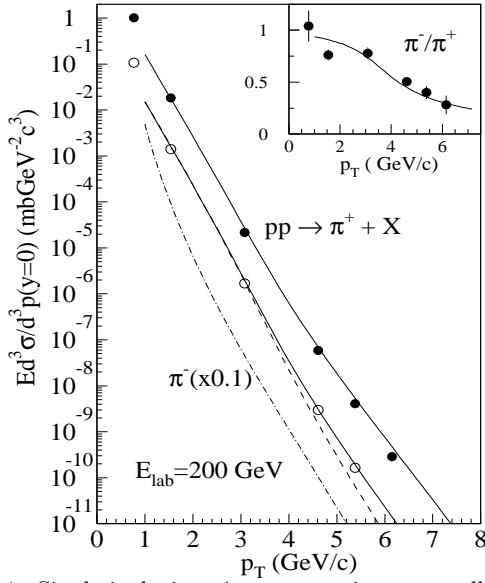


FIG. 1. Single-inclusive pion spectra in $p + p$ collisions at $E_{\text{lab}} = 200$ GeV. The solid lines are QCD parton model calculations with Q -dependent intrinsic k_T and the dot-dashed line is without. The dashed line is for a constant intrinsic $\langle k_T^2 \rangle_N = 1.5$ GeV². Experimental data are from Ref. [31]. The inserted figure shows the corresponding π^-/π^+ ratio.

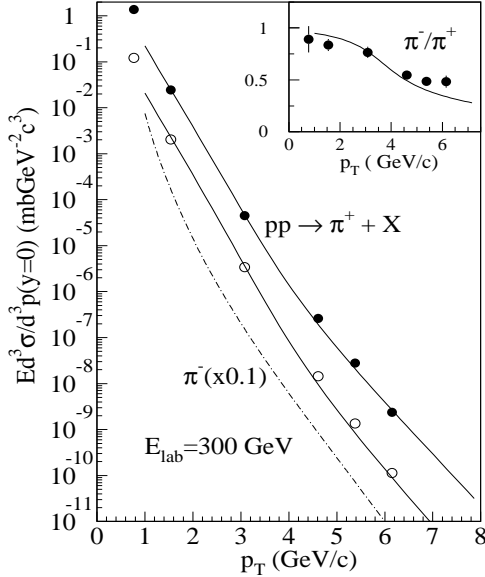


FIG. 2. Single-inclusive pion spectra in $p + p$ collisions at $E_{\text{lab}} = 300$ GeV. The solid lines are QCD parton model calculations with Q -dependent intrinsic k_T and the dot-dashed line is without. Experimental data are from Ref. [31]. The inserted figure shows the corresponding π^-/π^+ ratio.

In Fig. 1, we also show as a dashed line the calculated spectrum of π^- with a constant average initial transverse momentum, $\langle k_T^2 \rangle_N = 1.5$ GeV². At large p_T such a constant intrinsic $\langle k_T^2 \rangle_N$ underestimates the experimental data. This demonstrates phenomenologically why the Q dependence of $\langle k_T^2 \rangle_N$ is needed to fit the experiment-

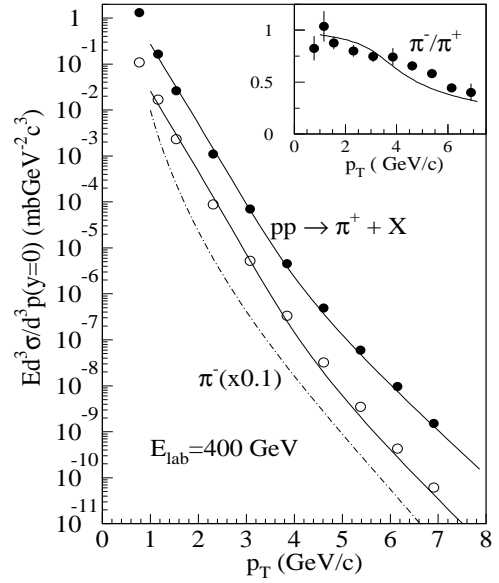


FIG. 3. The same as Fig. 2, except at $E_{\text{lab}} = 400$ GeV

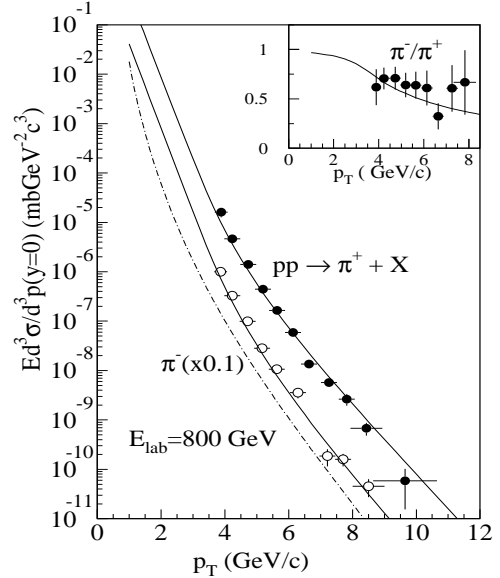


FIG. 4. The same as Fig. 2, except at $E_{\text{lab}} = 800$ GeV and the experimental data are from Ref. [32]

al data in particular at large values of $x_T = 2p_T/\sqrt{s}$ for the choice of Q^2 we used. One can also use an alternative choice for $Q^2 = P_T^{\text{jet}2} = (p_T/z_c)^2$. In this case we found that one does not need Q -dependent $\langle k_T^2 \rangle_N$ to fit the data in pp collisions. However, one needs to use an energy-dependent initial parton transverse momentum distribution. We also checked that as energy increases such a Q -dependent $\langle k_T^2 \rangle_N$ is not needed anymore to fit the experimental data. However, the available experimental data at collider energies are only limited to finite p_T range where $x_T = 2p_T/\sqrt{s}$ is not so large as compared to the low energy data. In the following we will use the Q -dependent initial momentum distribution

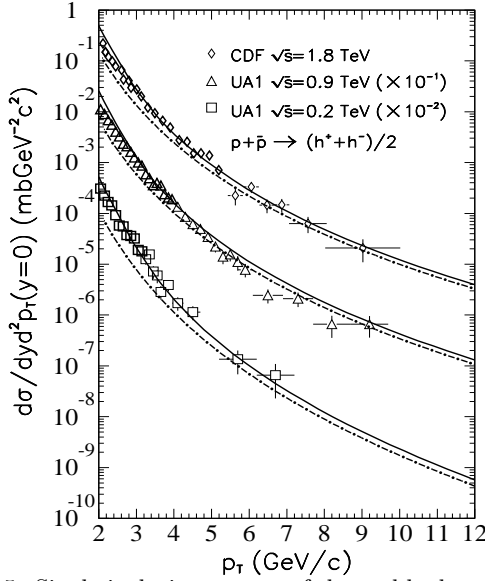


FIG. 5. Single-inclusive spectra of charged hadrons in $p + \bar{p}$ collisions at $\sqrt{s} = 200, 900, 1800$ GeV. The solid lines are QCD parton model calculations with intrinsic k_T and the dot-dashed lines are without. Experimental data are from Refs. [33,34].

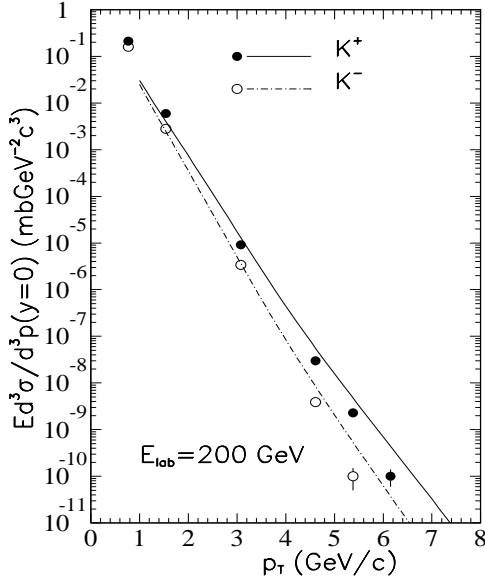


FIG. 6. Single-inclusive spectra of charged kaons in $p + p$ collisions at $E_{\text{lab}} = 200$ GeV. The lines are QCD parton model calculations with intrinsic k_T and data are from Ref. [31].

for our study of high p_T spectra in $p + A$ and $A + A$ collisions.

In the inserted boxes in Figs. 1-4, we also plot π^-/π^+ ratio as a function of p_T . At higher p_T , particle production is more dominated by the leading hadrons from valence quark scattering. Since there are more up-quarks than down-quarks in $p + p$ system, one should expect the ratio to become smaller than 1 and decrease with p_T . The QCD parton model calculations describe this isospin de-

pendence of the spectra very well. For other flavors of hadrons, we show in Figs. 5-7 the calculated kaon spectra at different energies. We see that the agreement with experimental data is very good, except at large p_T at low energies which may be improved by using better fragmentation functions for kaons. One can also notice that K^- spectra at large p_T are about a factor of 10 smaller than K^+ . This is because the content of strange quarks in a nucleon is much smaller than up (down) quarks which are responsible for leading K^+ (K^0) hadron production. We did not calculate and compare the spectra of produced protons and anti-protons because there is no very accurate parameterization of the corresponding fragmentation functions.

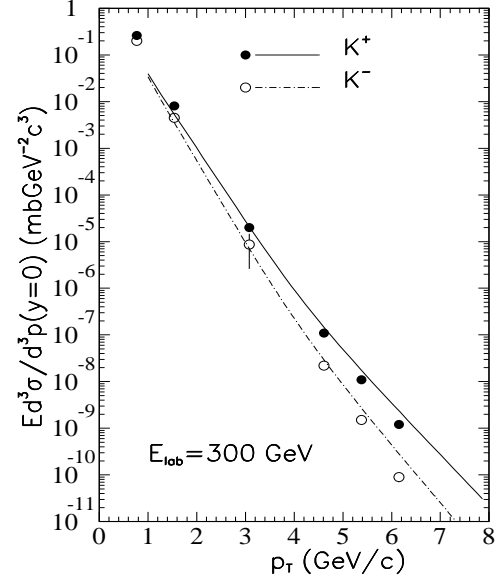


FIG. 7. The same as Fig. 6, except at $E_{\text{lab}} = 300$ GeV

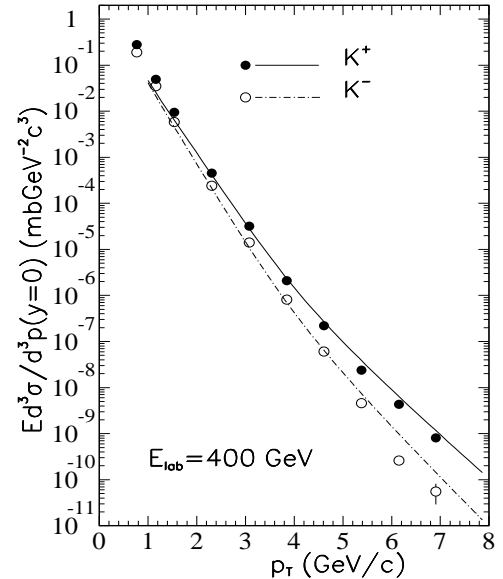


FIG. 8. The same as Fig. 6, except at $E_{\text{lab}} = 400$ GeV

III. $P + A$ COLLISIONS: K_T BROADENING

Since single-inclusive particle spectra at high p_T in $p + A$ collisions have been shown, both experimentally [31] and theoretically [35], to be sensitive to multiple initial-state scattering, or Cronin effect, it is important that we take into account this effect here in order to have a quantitative study of the change of particle spectra at high p_T in nucleus-nucleus collisions. One can study the Cronin effect in a model of multiple parton scattering [35]. In such a model, one uses Glauber multiple scattering formula to treat parton scattering between beam and target partons which is essential to incorporate the interference effect. For example, when contribution from double scattering is considered, one must also include the absorptive part of the single scattering which has a negative contribution. This absorptive part then cancels part of the contribution from double scattering. Consequently, the enhancement of particle spectra because of double scattering decreases with p_T in the form of $1/p_T^2$ and in general with \sqrt{s} [36]. One can also find that the dominant contribution in double scattering comes from the case where one of the scatterings is soft and large part of the final p_T of the produced jet comes from just one hard scattering. At much lower p_T the absorptive correction is much larger than the double scattering contribution and the spectra there is even suppressed. So the integrated cross section or average particle multiplicity at moderate p_T are not affected by multiple parton scattering even though the spectra are modified, as observed in the DY case [37]. This will provide us a justification for our following phenomenological treatment of multiple parton scattering in $p + A$ and $A + A$ collisions.

In this paper we assume that the inclusive differential cross section for large p_T particle production is still given by a single hard parton-parton scattering. However, due to multiple parton scattering prior to the hard processes, we consider the initial transverse momentum k_T of the beam partons is broadened. Assuming that each scattering provide a k_T kick which also has a Gaussian distribution, we can in effect just change the width of the initial k_T distribution. Then the single inclusive particle cross section in minimum-biased $p + A$ collisions is,

$$\begin{aligned} \frac{d\sigma_{pA}^h}{dyd^2p_T} &= K \sum_{abcd} \int d^2bt_A(b) \int dx_a dx_b d^2k_{aT} d^2k_{bT} \\ &\times g_A(k_{aT}, Q^2, b) f_{a/p}(x_a, Q^2) \\ &\times g_p(k_{bT}, Q^2) f_{b/A}(x_b, Q^2, b) \\ &\times \frac{D_{h/c}^0(z_c, Q^2)}{\pi z_c} \frac{d\sigma}{dt}(ab \rightarrow cd), \end{aligned} \quad (7)$$

where $t_A(b)$ is the nuclear thickness function normalized to $\int d^2bt_A(b) = A$. We will use Wood-Saxon form of nuclear distribution for $t_A(b)$ throughout this paper unless specified otherwise. The parton distribution per nu-

cleon inside the nucleus (with atomic mass number A and charge number Z) at an impact parameter b ,

$$\begin{aligned} f_{a/A}(x, Q^2, b) &= S_{a/A}(x, b) \left[\frac{Z}{A} f_{a/p}(x, Q^2) \right. \\ &\quad \left. + (1 - \frac{Z}{A}) f_{a/n}(x, Q^2) \right], \end{aligned} \quad (8)$$

is assumed to be factorizable into the parton distribution in a nucleon $f_{a/N}(x, Q^2)$ and the nuclear modification factor $S_{a/A}(x, b)$ which we take the parameterization used in HIJING [38] for now. The initial parton transverse momentum distribution inside a projectile nucleon going through the target nucleon at an impact parameter b is then,

$$g_A(k_T, Q^2) = \frac{1}{\pi \langle k_T^2 \rangle_A} e^{-k_T^2 / \langle k_T^2 \rangle_A}, \quad (9)$$

with a broadened variance

$$\langle k_T^2 \rangle_A(Q^2) = \langle k_T^2 \rangle_N(Q^2) + \delta^2(Q^2)(\nu_A(b) - 1). \quad (10)$$

The broadening is assumed to be proportional to the number of scattering $\nu_A(b)$ the projectile suffers inside the nucleus. For the purpose of considering the impact-parameter dependence of the k_T broadening, we will simply assume a hard sphere nuclear distribution. Therefore,

$$\nu_A(b) = \sigma_{NN} t_A(b) = \sigma_{NN} \frac{3A}{2\pi R_A^2} \sqrt{1 - b^2/R_A^2}, \quad (11)$$

where $R_A = 1.12A^{1/3}$ fm and σ_{NN} is the inelastic nucleon-nucleon cross section.

We also assume that k_T broadening during each nucleon-nucleon collision δ^2 also depends on the hard momentum scale $Q = P_T^{\text{jet}}$, the transverse momentum of the produced parton jet. Such a dependence is easy to understand by comparing the role of P_T^{jet} in the multiple parton scattering case to that in the initial and final state radiation associated with a hard jet production. For example, unless the multiple scatterings are far separated in space-time such that the propagating parton becomes on-shell, the interaction cross of the propagating parton will depend on its virtuality which in return depends on the scale of the hard scattering. In Ref. [39], the k_T broadening in a multiple parton scattering scenario is related to the ratio of the leading-twist parton distributions and higher-twist multiple parton distributions which apparently depends on the momentum scale of the hard parton scattering. In this paper, we will use the following scale-dependent k_T broadening per nucleon-nucleon collision,

$$\delta^2(Q^2) = 0.225 \frac{\ln^2(Q/\text{GeV})}{1 + \ln(Q/\text{GeV})} \text{ GeV}^2/c^2. \quad (12)$$

Such a functional form is chosen to best fit the experimental data in $p + A$ collisions. For $Q = 2 \sim 3$ GeV,

$\delta^2 = 0.064 \sim 0.129 \text{ GeV}^2/c^2$, which is consistent with the value obtained from the analysis of k_T broadening for J/Ψ production in $p+A$ collisions [40,41]. We should point out that this value can be different from what one gets from the analysis of k_T broadening of Dell-Yan lepton pairs in $p+A$ collisions, where mainly quarks and anti-quarks are involved in the hard processes and there is no collision effect after the Dell-Yan production point.

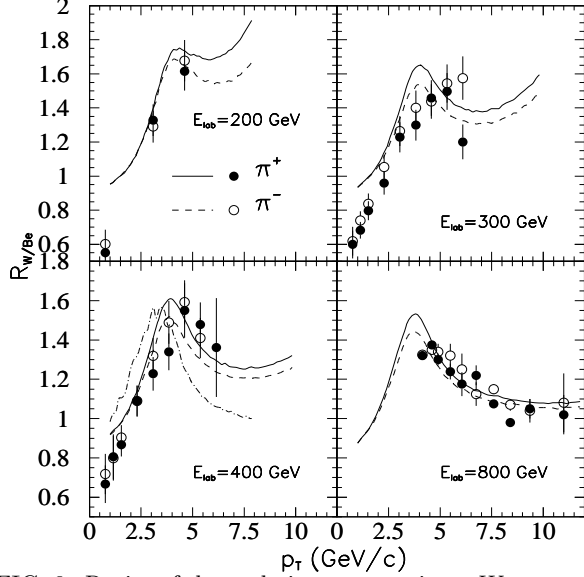


FIG. 9. Ratios of charged pions spectra in $p+W$ over $p+Be$ each normalized by the atomic number of the target nucleus. The lines are the parton model calculation with k_T broadening due to multiple parton scattering. The dot-dashed line in the third panel is the calculated result with a constant k_T broadening $\delta^2 = 0.1 \text{ GeV}^2$. Experimental data are from Refs. [31,42].

Using the above assumptions, we calculate the large p_T particle spectra in $p+A$ collisions and compare to the spectra in $p+p$ collisions. As we show in the last Section, particle spectra has a strong isospin dependence in the parton model. In order to minimize this known isospin dependence of the spectra in our study of nuclear dependence of the spectra, we compare the spectra of $p+A$ for heavy target with that for a very light nuclear target. Shown in Fig. 9 are our calculated ratios of charged pion spectra (solid lines for π^+ and dot-dashed lines for π^-) in $p+W$ over that of $p+Be$ each normalized by the atomic number of the target nucleus. The experimental data are from Ref. [31,42]. If there was no nuclear dependence due to multiple scattering, the ratios would have a flat value of 1, modulo the residue isospin dependence because of the small isospin asymmetry of the target nucleus. The difference between our calculated ratios for π^+ and π^- gives the order of this isospin asymmetry effect. One can get rid of this effect by using the ratios of $\pi^+ + \pi^-$ spectra. As shown in the figure, our model can roughly describe the general feature of the nuclear dependence of the spectra at large p_T due to multiple parton

scattering. The ratios should become smaller than 1 at very small p_T because of the absorptive processes as we have mentioned. But here our perturbative calculation will eventually breaks down because of the small momentum scale. At larger p_T , the spectra are enhanced because of multiple parton scattering. As p_T increases further, the ratios decrease again and saturate at about 1. The decrease follows the form of $1/p_T^2$ consistent with the general features of high twist processes. Since the transverse momentum broadening due to multiple parton scattering is finite, its effect will eventually become smaller and disappear. Therefore, the p_T location of the maximum enhancement can give us the scale of average transverse momentum broadening. We also show in the third panel as dot-dashed line our calculation with a constant k_T broadening, $\delta^2 = 0.1 \text{ GeV}^2$. While the result has the general feature of the experimental data, the fit is not as good as the scale-dependent k_T broadening.

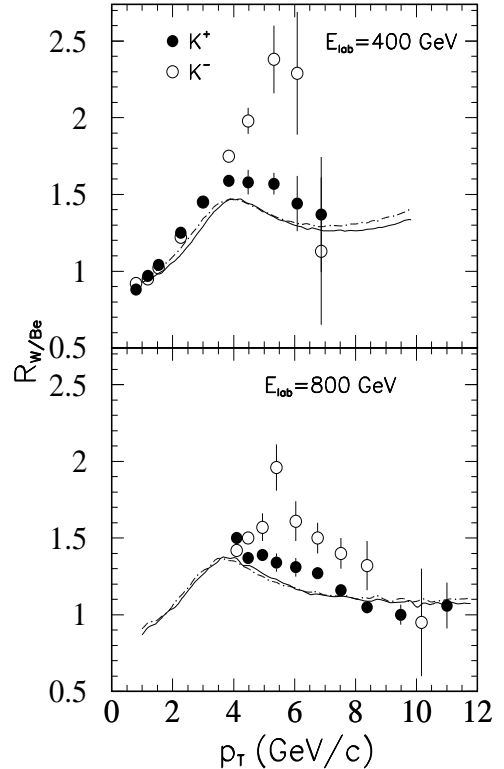


FIG. 10. Ratios of charged kaons spectra in $p+W$ over $p+Be$ each normalized by the atomic number of the target nucleus. The solid (dot-dashed) lines are the parton model calculation for K^+ (K^-) with k_T broadening due to multiple parton scattering. Experimental data are from Refs. [31,42].

At low energies, $E_{\text{lab}} = 200 \text{ GeV}$ for example, the full structure of multiple scattering cannot be revealed because of the finite phase space constrained by the kinetic limit. Furthermore, the ratio will increase at large $x_T = 2p_T/\sqrt{s}$ because of the so-called EMC effect [43] on nuclear structure function at large x which is incorporated into the parameterization of the nuclear modification factor for parton distributions inside a nucleus. One

word of caution should also be given here. Our calculation might not be reliable anymore near the boundary of the available phase space ($p_T \sim \sqrt{s}/2$) where overall energy and momentum conservation is very important which is not imposed in QCD parton model.

We have also calculated the kaon spectra in pA collisions. Show in Fig. 10 are the ratios of the kaon spectra in $p+W$ over $p+Be$ collisions (normalized by the atomic number). The agreement with the experimental data is quite good for K^+ . However, for K^- there is quite sizable discrepancy between the calculation and the data. This might be due to the flavor dependence of the k_T broadening suffered by the quarks. It is quite possible that \bar{u} quarks have larger scattering cross sections as they propagate through the nucleus and therefore experience larger k_T broadening than \bar{s} quarks. Such effects will result in the observed difference in nuclear enhancement between K^- and K^+ .

IV. $A + A$ COLLISIONS: JET QUENCHING ?

It is straightforward to incorporate the initial k_T broadening due to multiple parton scattering in $A + A$ collisions. Since partons from projectile and target beam both suffer multiple scattering before the hard process, their initial transverse momentum distributions are broadened as given by Eq. (9). In addition, the produced parton jets have to propagate through a dense medium and interact with other produced particles in the medium. Theoretical studies [3–5] show that a fast parton propagating inside a dense partonic medium will suffer considerable amount of energy loss. Since large p_T particles are produced through the fragmentation of parton jets with large transverse energy, parton energy loss will definitely lead to the suppression of large p_T particles. One can describe this suppression phenomenologically via effective fragmentation functions which are modified from their original forms in the vacuum [7–9]. In this approach one can then calculate the single inclusive particle spectra at large p_T in $A + B$ collisions as,

$$\begin{aligned} \frac{d\sigma_{AB}^h}{dyd^2p_T} &= K \sum_{abcd} \int d^2b d^2r t_A(r) t_B(|\mathbf{b} - \mathbf{r}|) \\ &\times \int dx_a dx_b d^2k_{aT} d^2k_{bT} \\ &\times g_A(k_{aT}, Q^2, r) g_B(k_{bT}, Q^2, |\mathbf{b} - \mathbf{r}|) \\ &\times f_{a/A}(x_a, Q^2, r) f_{b/B}(x_b, Q^2, |\mathbf{b} - \mathbf{r}|) \\ &\times \frac{D_{h/c}(z_c, Q^2, \Delta L)}{\pi z_c} \frac{d\sigma}{dt}(ab \rightarrow cd), \end{aligned} \quad (13)$$

where $D_{h/c}(z_c, Q^2, \Delta L)$ is the modified effective fragmentation function for produced parton c which has to travel an average distance ΔL inside a dense medium. We will not elaborate on the modeling of the modified fragmentation functions [7–9] here, except pointing out that it

depends on two parameters: the energy loss per scattering ϵ_c and the mean free path λ_c for a propagating parton c . The energy loss per unit length of distance is then $dE_c/dx = \epsilon_c/\lambda_c$. We also assume that a gluon's mean free path is half of a quark and then the energy loss dE/dx is twice that of a quark. In principle, the energy loss dE/dx should depend on local parton density or temperature, the parton's initial energy, the total distance and whether there is expansion [44]. These possible features will influence the final hadron spectra and their phenomenological consequences have been studied in detail in a previous publication [9]. In this paper, we will simply assume a constant energy loss and study the average effect of parton energy loss in dense matter.

After integration over the impact parameter space, Eq. (13) will give us the inclusive hadron spectra for minimum-biased events of $A+B$ collisions. If one neglects nuclear effects (parton shadowing, k_T broadening and jet quenching), the resultant spectra should be exactly proportional to AB which is the averaged number of binary nucleon-nucleon collisions. Very often experimentalists also measure inclusive hadron spectra for certain classes of events with different centrality cuts (total transverse energy E_T or charged multiplicity in the central region). Since one cannot directly measure the impact parameters in each class of centrality, a theoretical model of correlation between impact parameter b and the total transverse energy E_T (or charged multiplicity) has to be introduced in order to calculate the inclusive cross section for events with different centrality cuts. In this paper, we will use the correlation function $P_{AB}(E_T, b)$ [normalized to $\int dE_T P_{AB}(E_T, b) = 1$] introduced in Ref. [45] which assumes a wounded nucleon model for transverse energy production. The differential cross section

$$\frac{d\sigma_{AB}}{dE_T} = \int d^2b [1 - e^{-\sigma_{NN} T_{AB}(b)}] P_{AB}(E_T, b) \quad (14)$$

has been shown to reproduce the experimental data of NA35 and NA49 [46] very well, where $T_{AB}(b) = \int d^2r t_A(\mathbf{r}) t_B(\mathbf{b} - \mathbf{r})$ is the overlap function for $A + B$ collisions at impact parameter b . We refer readers to Ref. [45] for the details of this model.

After incorporating the correlation function between impact parameter and transverse energy in Eq. (13), the inclusive cross section of large p_T hadron production in events with centrality cut $E_T \in (E_T^{\min}, E_T^{\max})$ can be shown to be roughly proportional to

$$\langle N_{\text{binary}} \rangle \equiv \int d^2b \int_{E_T^{\min}}^{E_T^{\max}} dE_T P_{AB}(E_T, b) T_{AB}(b) \quad (15)$$

which is just the average number of binary nucleon-nucleon collisions. For minimum-biased events, this number is just AB . Since the actual $d\sigma_{AB}/dE_T$ distribution depends on each experiment's coverage of phase space and the detector's calibration, we will choose the values of E_T^{\min} and E_T^{\max} so that the fraction of the integrated

cross section in Eq. (14) within the E_T range matches the experimental value of a given centrality cut.

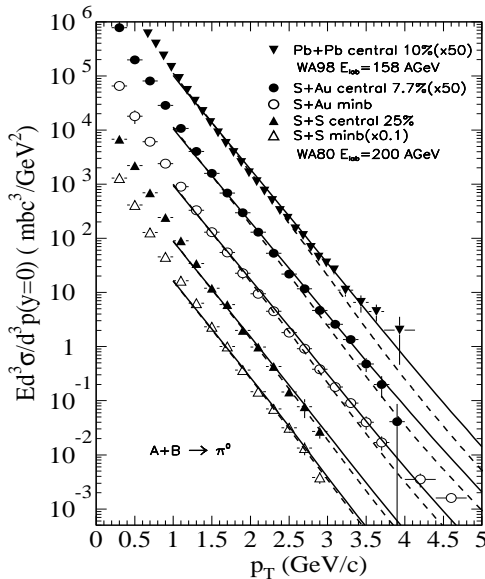


FIG. 11. Single-inclusive spectra of π^0 in $S+S$, $S+Au$ and $Pb+Pb$ collisions (both minimum-biased and central events) at the CERN SPS energies. The solid lines are QCD parton model calculations with k_T broadening due to initial multiple scattering and the dashed lines are without. Experimental data are from Refs. [10,11].

Shown in Fig. 11 are the calculated inclusive spectra for produced π_0 in $S+S$, $S+Au$ and $Pb+Pb$ collisions, both minimum-biased and central events. The QCD parton model calculations with the k_T broadening due to initial multiple scattering (solid lines) agree with the experimental data (WA80 and WA98) [10,11] well very at p_T above 1 GeV/c. No parton energy loss has been assumed in the calculations. The dashed lines are the spectra in pp collisions at the same energy multiplied by the nuclear geometrical factor as given in Eq. (15). The difference between the solid and dashed lines is simply caused by effects of k_T broadening and nuclear modification of parton distributions inside nuclei. These effects are similar as in $p+A$ collisions and are more important in collisions of heavier nuclei. Without these nuclear effects the high p_T hadron spectra in $A+B$ collisions are exactly proportional to the average number of binary collisions as shown by the dashed lines. This is a common characteristic of hard processes in $p+A$ and $A+B$ collisions. Because of absorptive processes, low p_T particle production, which can be considered as coherent over the dimension of nuclear size, has much weaker A -dependence. In the wounded-nucleon model, soft particle production cross section is proportional to the average number of wounded nucleons which is much smaller than the number of binary collisions. By studying the transition of the scaling property of the hadron spectra from low to high p_T values, one can then determine at what p_T value the underlying mechanism of hadron production become

dominated by hard processes. This will be the subject of another study.

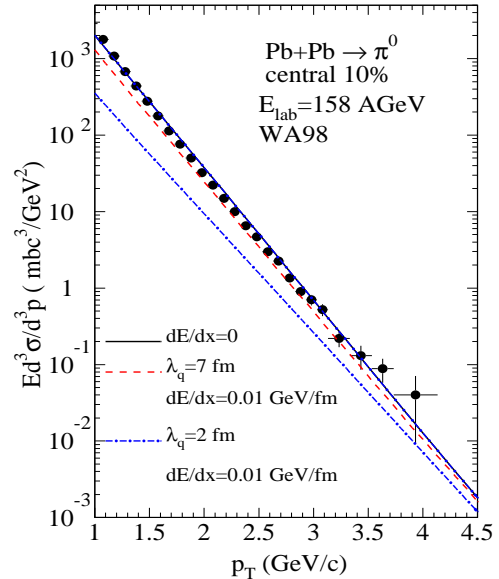


FIG. 12. Parton model calculations of single-inclusive spectra of π^0 with different values of parton energy loss dE_q/dx and mean free path λ_q in central $Pb+Pb$ collisions at $E_{lab} = 158$ AGeV/. The Experimental data are from Ref. [11].

In the parton model calculations presented so far, effects of parton energy loss have not been considered yet. If there is parton energy loss and the radiated gluons become incoherent from the leading parton, the resultant leading hadron spectra at large p_T from the parton fragmentation should be suppressed as compared to $p+p$ and $p+A$ collisions. As we have shown, however, the parton model calculations without parton energy loss fit the experimental data very well. To find out the consequences of an effective parton energy loss in the hadron spectra and how the experimental data constrain its value at the CERN SPS energy, we show in Fig. 12 as the dot-dashed line the calculated π_0 spectra in central $Pb+Pb$ collisions with $dE/dx = 0.01$ GeV/fm. In the calculation (dot-dashed line) we have assumed the energy loss per scattering $\epsilon_q = 0.02$ GeV and the mean free path $\lambda_q = 2$ fm ($dE_q/dx = \epsilon_q/\lambda_q$). In our model [7,8], the modification of the fragmentation function is sensitive to both parameters. The mean free path is a measure of the strength of the interaction between the leading parton and medium while energy loss dE/dx reflects the degree of attenuation induced by the interaction. Shown as the dashed line in Fig. 12 is a calculation with the same $dE_q/dx = 0.01$ GeV/fm but with a mean free path $\lambda_q = 7$ fm which is about the average total length a parton will travel through in a cylindrical system with a radius of a Pb nucleus. Such a scenario of weak interaction and small energy loss is barely consistent with the systematics of the experimental data. It is clear that the large p_T spectra in $Pb+Pb$ collisions at the CERN SPS energy

can put very stringent limits on the interaction of energetic partons with dense medium and the induced energy loss. Within the parton model, one can exclude from the observed hadron spectra a parton energy loss larger than $dE_q/dx = 0.01$ GeV/fm and a mean free path shorter than $\lambda_q = 7$ fm. This is much smaller than the most conservative estimate of parton energy loss in a dense medium [3–5]. According to the most conservative estimate in Ref. [5], a quark should still have an energy loss $dE/dx \approx 0.2$ GeV/fm in a cold nuclear matter of 10 fm in transverse size. Our constraint is 20 times smaller than this conservative limit, even though the dense matter in AA collisions should at least consist of hot hadronic matter which is much denser than the normal cold nuclear matter.

There are several implications one can draw from this analysis. Most of the recent theoretical estimates of parton energy loss are based on a scenario of a static and infinitely large dense parton gas. If the system produced in a central $Pb + Pb$ collision only exists for a period of time shorter than the interaction mean free path of the propagating parton, one then should not expect to see any significant parton energy loss. Using the measured transverse energy production $dE_T/d\eta \approx 405$ GeV [46] and a Bjorken scaling picture, one can indeed estimate [12] that the life time of the dense system in central $Pb + Pb$ collisions is only about $2 - 3$ fm/c before the density drops below a critical value of $\epsilon_c \approx 1$ GeV/fm³. Even if we assume that a dense partonic system is formed in central $Pb + Pb$ collisions, this optimistic estimate of the life time of the system could still be smaller than the mean free path of the propagating parton inside the medium. Thus, one does not have to expect a significant effect of parton energy loss on the final hadron spectra at large p_T . Otherwise, it will be difficult to reconcile the absence of parton energy loss with the strong parton interaction which maintains a long-lived partonic system.

Another conclusion one can also make is that the dense hadronic matter which has existed for a period of time in the final stage of heavy-ion collisions does not cause any apparent parton energy loss or jet quenching. A high p_T physical pion from jet fragmentation has a very long formation time. One does not have to worry about its scattering with other soft hadrons in the system which could cause suppression of high p_T pion spectra. We still do not understand the reason why a fragmenting parton does not lose much energy when it propagates through a dense hadronic matter. However, it might be related to the absence of energy loss to the quarks and anti-quarks prior to Drell-Yan hard processes in $p + A$ and $A + A$ collisions. This observation will make jet quenching a better probe of a long-lived partonic matter since one does not have to worry about the complications arising from the hadronic phase of the evolution. If one observes a dramatic suppression of high p_T hadron spectra at the BNL RHIC energy as predicted [2,7–9], then it will clearly indicate an initial condition very different from what has been reached at the CERN SPS energy.

V. PREDICTIONS AT RHIC: IMPORTANCE OF $P + A$ EXPERIMENTS

As we have seen in the calculation of large p_T particle spectra in $p + p(\bar{p})$ collisions, the initial k_T smearing is most important at low energies where the original particle spectra from jet production falls off very rapidly with p_T . The effect of k_T broadening due to multiple parton scattering in $p + A$ collisions follows the same line. At higher energies ($\sqrt{s} > 50$ GeV), when the original jet spectra become much flatter, the effects of both the initial or intrinsic k_T and nuclear k_T broadening will become smaller, as we can clearly see in both our calculation and the experimental data in Figs. 8 and 5. At the BNL RHIC collider energy, the nuclear enhancement of large p_T spectra will then become smaller. Shown in Fig. 13 are the calculated ratios of charged hadron spectra in $p + Au$ over that in $p + p$ normalized to the averaged number of binary nucleon collisions,

$$R_{AB}(p_T) \equiv \frac{d\sigma_{AB}^h/dy d^2p_T}{\langle N_{\text{binary}} \rangle d\sigma_{pp}^h/dy d^2p_T} \quad (16)$$

where $\langle N_{\text{binary}} \rangle = A$ for minimum-biased events of $p + A$ collisions. We can see that the enhancement due to multiple parton scattering will not disappear at the RHIC energy. There is still about 20–50% enhancement at moderately large p_T around 4 GeV/c. It then disappears very quickly at larger p_T .

At the BNL RHIC energy, $\sqrt{s} = 200$ GeV, nuclear modification of the gluon distribution will become important for hadron spectra at moderately large p_T . Such modifications for quark distributions have been measured in deeply inelastic lepton-nucleus collisions [47]. However, the nuclear effects on the gluon distribution have not been directly measured. In the calculation shown as the dashed line in Fig. 13 we have used a recent parameterization of the nuclear modification factors $S_{a/A}(x, Q^2)$ [Eq. (8)] by Eskola, Kolhinen and Salgado [47] which is based on global fits to the most recent collection of data available and some model on nuclear modification of gluon distribution. This result is quite different from the calculation using HIJING [38] parameterization. In EKS98 parameterization, QCD evolution equation has been used to take into account the Q^2 scale dependence of the nuclear modification which is absent in HIJING parameterization. Because of QCD evolution, the nuclear shadowing effects at small x become smaller with increasing Q^2 . This is why the hadron spectrum in $p + A$ collisions using EKS parameterization is larger than that using HIJING parameterization. The EKS98 parameterization also has a gluon anti-shadowing which is larger than any parameterizations before. Hadron spectra in a moderately large p_T range at the BNL RHIC energy mainly come from fragmentation of gluon jets. This is

why the EKS98 result in Fig. 13 is larger than that without nuclear modification of parton distributions (solid line) in the p_T range where anti-shadowing becomes relevant. In addition, the impact parameter dependence of the nuclear modification of parton distributions is not implemented in the calculation using EKS98 parameterization. Though EKS98 parameterization has taken into account information about nuclear shadowing of gluon distribution from the measured scale evolution of the structure functions of different nuclei, such a procedure is also model dependent. The result will depend on whether one includes the higher gluon-density terms in the QCD evolution equation. Before a direct measurement of the nuclear effect on gluon distribution inside a nucleus, one should consider it one of the uncertainties in predicting the hadron spectra in $p + A$ collisions at the BNL RHIC energy.

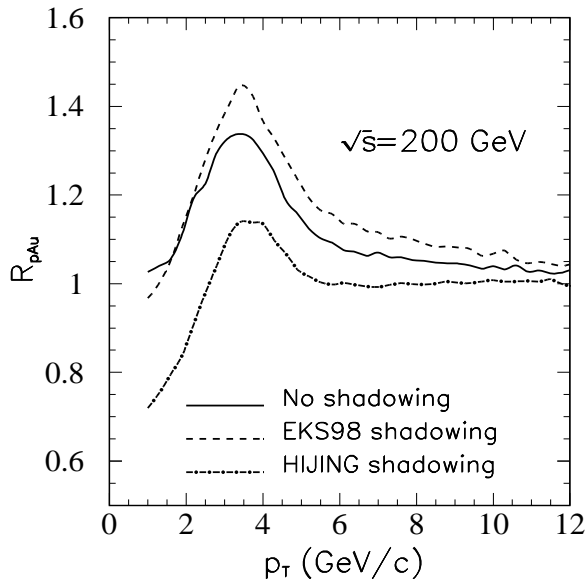


FIG. 13. Predictions for the ratio of single charged hadron spectra in $p + Au$ over $p + p$ collisions normalized by the average number of binary collisions (or A) at $\sqrt{s} = 200$ GeV. Different lines are for different parameterizations of shadowing or nuclear modification of parton distributions.

As pointed out in Sec. II, the shape of the nuclear modification of hadron spectra at large p_T is a result of multiple parton scattering inside a nucleus and their corresponding absorptive corrections. The shape is remarkably similar to the nuclear modification of parton distributions in small x region (shadowing and anti-shadowing) can be explained in a multiple scattering model [48]. Therefore, one could have double counted the same effect if he includes both the Cronin effect (or k_T broadening) and nuclear modification of the parton distributions in the calculation of hadron spectra in $p + A$ collisions. Given these uncertainties, it is therefore extremely important to have a systematic study of $p + A$ collisions at the BNL RHIC energy. Such an effort is pivotal to un-

ravel any other effects in the hadron spectra caused by the formation of dense partonic matter in $A + A$ collisions.

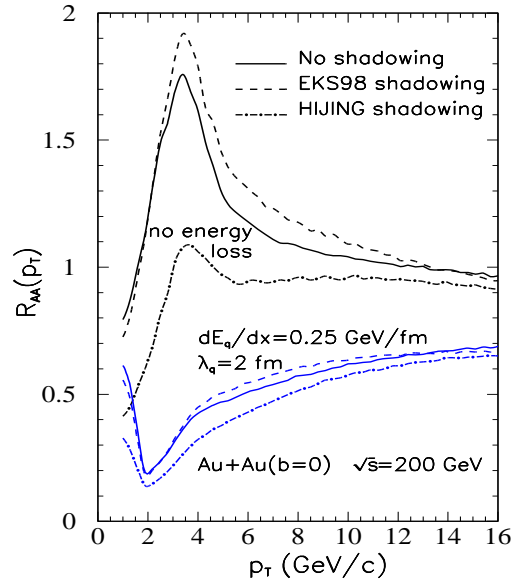


FIG. 14. Predictions for the ratio of single charged hadron spectra in central $Au + Au$ over $p + p$ collisions normalized by the average number of binary collisions at $\sqrt{s} = 200$ GeV. Different lines are for different parameterizations of shadowing or nuclear modification of parton distributions. The upper set of lines are without parton energy loss and lower set are with parton energy loss $dE_q/dx = 0.25$ GeV/fm and mean free path $\lambda_q = 2$ fm.

Shown in Fig. 14 are predictions for the ratio of hadron spectra in central $Au + Au$ collisions over that in $p + p$ [as defined in Eq. (16)], with and without parton energy loss. In both cases, uncertainties in the nuclear modification of parton distributions are still important. If there is significant parton energy loss in the dense medium at the BNL RHIC energy, leading hadrons from jets which are produced in the center of a overlapped region will be suppressed. As we can see this is where the Cronin effect and the effect of nuclear modification of parton distributions are the strongest. This is why the relative uncertainty caused by nuclear parton distributions is reduced compared to the case when there is no parton energy loss. In the calculations, we also included possible contribution to hadron production from soft processes at low p_T as implemented in Ref. [9]. When hadrons from hard parton scattering are suppressed, these soft contributions might become dominant at intermediate p_T . This is why the spectra ratio increases again when $p_T < 2$ GeV/c in Fig. 14. The spectra at lower p_T could also be broadened due to the rescattering effect which also drives the system into local equilibrium. We should emphasize that the soft particle spectra we use here are extremely schematic and qualitative. This is another uncertainty one should keep in mind which could affect the shape of spectra at intermediate p_T . Nevertheless, the spectra at moderately large p_T is most sensitive to parton energy loss in dense

medium. Therefore, one should study hadron spectra in $p + A$ very carefully and pin down the uncertainty due to multiple scattering and parton (anti)shadowing as much as possible. Only then one can draw more accurate conclusions about parton energy loss from the single inclusive hadron spectra in heavy-ion collisions at the BNL RHIC energy.

VI. CONCLUSIONS AND DISCUSSIONS

In this paper, we have analyzed systematically large p_T hadron spectra in $p + p$, $p + A$ and $A + A$ collisions from CERN SPS to BNL RHIC energies within a QCD parton model. We found that both the initial k_T in $p + p$ collisions and the k_T broadening due to multiple parton scattering in $p + A$ collisions are important to describe the experimental data within the parton model calculations. The value of initial k_T in order to fit the data is found to be larger than the conventional value of 300 – 500 MeV for intrinsic k_T according to the uncertainty principle. This finding is also consistent with analysis of Drell-Yan data [26] and recent study [20] of direct photon and pion production at around CERN SPS energy range. In both studies, one found that an intrinsic k_T of the order of 1 GeV/ c is needed to describe the data within NLO parton model calculations. Since we only used LO pQCD calculation, we have to introduce some Q^2 dependence of the initial k_T induced by initial-state radiation processes in order to fit the experimental data at low collider energies. This parton model with phenomenologically included intrinsic k_T describes very well the energy and isospin dependence of the hadron spectra.

The parton model calculation can also describe the large p_T pion spectra in heavy-ion collisions at the CERN SPS energies very well, both the A or centrality and energy dependence. There is no evidence of proposed parton energy loss caused by dense partonic matter. Based on recent theoretical estimates [3–5] of parton energy loss in dense partonic matter, one should expect a parton energy loss in the order of $dE/dx \sim 2 - 4$ GeV/fm. The absence of such energy loss in large p_T hadron spectra implies that either there is no such dense partonic matter formed or the life time of such medium is smaller than the mean free path of the parton interaction inside such a medium. It also tells us that the hadronic matter which must have existed for a period of time in heavy-ion collisions at the CERN SPS will not cause apparent energy loss or jet quenching effect. Therefore, if one observes suppression of high p_T hadrons at the BNL RHIC energy, it will unambiguously reflect an initial condition very different from what has been achieved at the CERN SPS.

As pointed out in a recent paper by Gyulassy and Levai [49], even though HIJING [38] Monte Carlo model fails to reproduce the hadron spectra in $p + p$, $p + A$ and $S + S$ (and $S + Au$) data at the CERN SPS energy, it acci-

dentally reproduces the central $Pb + Pb$ data very well because some unique form of transverse momentum kick introduced in HIJING to the end-points of strings each time they suffer an interaction. A hydrodynamic model calculation [50] with significant transverse expansion can also be used to fit the high p_T hadron spectra. So one should wonder how one can make sure that the high p_T hadron production in $A + B$ collisions is indeed dominated by hard parton scattering. As we have proposed earlier [12,51], measurement of two-particle correlation in azimuthal angle in the transverse plane should be able to distinguish these models from parton model. In the parton model, jets are always produced in pairs and back-to-back in the transverse plane. High p_T particles from jet fragmentation should then have strong back-to-back correlation. Neither hydro-dynamical model nor multiple scattering of string end-points can give such correlation. If such correlation is seen, one can then study the p_T dependence of the correlation to find out at what p_T value the correlation disappears. One can then at least quantify above what p_T value thermal-hydro model can be ruled out as the underlying particle production mechanism. Such a study is underway and will be reported elsewhere.

At the BNL RHIC energy, things will become a little cleaner. As pointed out by Gyulassy and Levai [49], most of hadron production at moderate p_T will be dominated by gluon production and are not influenced by the fate of strings or valence quarks (except net baryons). The spectra should be more sensitive to parton energy loss. Moreover, the hadron spectra in $p + A$ and $A + A$ collisions without parton energy loss will have a unique shape similar to those shown in Fig. 13 at high energies which can not be fully revealed at the CERN SPS energies. This unique shape will be hard to explain by any thermal-hydrodynamic models. However, as we have demonstrated in the previous section, there are still some uncertainties related to nuclear modification of parton distributions. It is therefore very important to have a systematic study of $p + p$, $p + A$ and $A + A$ collisions at the BNL RHIC energy in order to make more quantitative conclusions about parton energy loss in dense medium from the single hadron spectra at moderately high p_T .

ACKNOWLEDGEMENTS

This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Divisions of Nuclear Physics, of the U.S. Department of Energy under Contract No. DE-AC03-76SF00098 and DE-FG03-93ER40792. The author wish to thank K. J. Eskola for comments and for providing the EKS98 parameterization of the nuclear modification of parton distributions.

-
- [1] M. Gyulassy and M. Plümer, Phys. Lett. **B243**, 432 (1990).
- [2] X.-N. Wang and M. Gyulassy, Phys. Rev. Lett. **68**, 1480 (1992).
- [3] M. Gyulassy and X.-N. Wang, Nucl. Phys. **B420**, 583 (1994); X.-N. Wang, M. Gyulassy and M. Plümer, Phys. Rev. **D 51**, 3436 (1995).
- [4] R. Baier, Yu. L. Dokshitzer, S. Peigné and D. Schiff, Phys. Lett. **B345**, 277 (1995).
- [5] R. Baier, Yu. L. Dokshitzer, A. Mueller, S. Peigné and D. Schiff, Nucl. Phys. **B484**, 265 (1997).
- [6] M. Gyulassy and M. Plümer, Nucl. Phys. **B346**, 1 (1990).
- [7] X.-N. Wang and Z. Huang, Phys. Rev. C **55**, 3047 (1997).
- [8] X.-N. Wang, Z. Huang and I. Sarcevic, Phys. Rev. Lett. **77**, 231 (1996).
- [9] X.-N. Wang, Phys. Rev. C **58**, 2321 (1998).
- [10] R. Albrecht *et al.*, WA80 Collaboration, Eur. Phys. J. C **5**, 255 (1998).
- [11] M. M. Aggarwal *et al.*, WA98 Collaboration, nucl-ex/9806004 (to be published).
- [12] X.-N. Wang, Phys. Rev. Lett. **81**, 2655 (1998).
- [13] For a review, see J. F. Owens, Rev. Mod. Phys. **59**, 465 (1987).
- [14] A. D. Martin, W. J. Stirling and R. G. Roberts, Phys. Lett. **B306**, 145 (1993); For a more recent global analysis see A. D. Martin, W. J. Stirling, R. G. Roberts and R. S. Thorne, DTP-98-10, March 1998, hep-ph/9803445.
- [15] J. Binnewies, B. A. Kniehl and G. Kramer, Z. Phys. **C65**, 471 (1995).
- [16] F. M. Borzumati and G. Kramer, Z. Phys. **C67**, 137 (1995); J. Binnewies, B. A. Kniehl and G. Kramer, Phys. Rev. **D53**, 3573 (1996).
- [17] D. C. Hom, *et al.*, Phys. Rev. Lett. **37**, 1374 (1976); D. M. Kaplan, *et al.*, Phys. Rev. Lett. **40**, 435 (1978).
- [18] BEATRICE Coll., M. Adamovich *et al.*, Phys. Lett. **B348**, 256 (1995); Phys. Lett. **B433**, 217 (1998).
- [19] M. N. Mangano, P. Nason and G. Ridolfi, Nucl. Phys. **B373**, 295 (1992).
- [20] L. Apanasevich, *et al.*, Phys. Rev. Lett. **81**, 2642 (1998); L. Apanasevich, *et al.*, Phys. Rev. D **59**, 074007 (1999).
- [21] G. Altarelli, R. K. Ellis, M. Greco and G. Martinelli, Nucl. Phys. **B246**, 12 (1984).
- [22] J. C. Collins, D. E. Soper and G. Sterman, Nucl. Phys. **B250**, 199 (1985).
- [23] Yu. L. Dokshitzer, D. I. Dyakonov and S. I. Troyan, Phys. Rep. **58**, 270 (1980).
- [24] J. C. Collin and R. K. Ellis, Nuc.; Phys. **B360**, 3 (1991); M. G. Ryskin, A. G. Shuvaev and Yu. M. Shabelski, hep-ph/9907507, 1999.
- [25] H. Li and G. Sterman, Nucl. Phys. **B381**, 129 (1992).
- [26] R. D. Field, *Applications of Perturbative QCD*, Frontiers in Physics Lecture, Vol. 77, Ch. 5.6 (Addison Wesley, 1989).
- [27] R. P. Feynman, R. D. Field, and G. C. Fox, Nucl. Phys. **B128**, 1 (1977); Phys. Rev. **D18**, 3320 (1978).
- [28] J. F. Owens and J. D. Kimel, Phys. Rev. **D18**, 3313 (1978).
- [29] Kaon fragmentation functions are adopted from, J. Binnewies, B. A. Kniehl and G. Kramer, Phys. Rev. **D52**, 4947 (1995).
- [30] K. J. Eskola and X.-N. Wang, Int. J. Mod. Phys. A **10**, 3071 (1995).
- [31] D. Antreasyan, *et al.*, Phys. Rev. **D19**, 764 (1979).
- [32] D. E. Jaffe, *et al.*, Phys. Rev. **D40**, 2777 (1989).
- [33] UA1 Collab., C. Albajar, *et al.*, Nucl. Phys. **B335**, 261 (1990).
- [34] F. Abe, *et al.*, Phys. Rev. Lett. **61**, 1819 (1988).
- [35] M. Lev and B. Petersson, Z. Phys. **C21**, 155 (1983); T. Ochiai, *et al.*, Prog. Theor. Phys. **75**, 288 (1986).
- [36] X.-N. Wang, Phys. Rep. **280**, 287 (1997).
- [37] P. Bordalo, *et al.*, NA10 Collaboration, Phys. Lett. **193B**, 373 (1987); D. M. Alde, E772 Collaboration, Phys. Rev. Lett. **66**, 2285 (1991).
- [38] X.-N. Wang and M. Gyulassy, Phys. Rev. D **44**, 3501 (1991); Comp. Phys. Comm. **83**, 307 (1994).
- [39] X. F. Guo, Phys. Rev. **D58**, 114033 (1998); 036001 (1998).
- [40] S. Gavin and M. Gyulassy, Phys. Lett. **B214**, 241 (1988).
- [41] D. Kharzeev, M. Nardi and H. Satz, Phys. Lett. **B405**, 14 (1997), and reference therein.
- [42] P. B. Straub, *et al.*, Phys. Rev. Lett. **68**, 452 (1992).
- [43] M. Arneodo *et al.*, EM Collaboration, Nucl. Phys. **B333**, 1 (1990).
- [44] R. Baier, Yu. L. Dokshitzer, A. H. Mueller and D. Schiff, hep-ph/9803473.
- [45] C. Lorenc, D. Kharzeev, M. Nardi and H. Satz, Z. Phys. **C 74**, 307 (1997).
- [46] T. Alber *et al.* (NA49 experiment), Phys. Rev. Lett. **75**, 3814 (1995).
- [47] K.J. Eskola, V.J. Kolhinen, C.A. Salgado, hep-ph/9807297; K.J. Eskola, V.J. Kolhinen, P.V. Ruuskanen, hep-ph/9802350; and references therein.
- [48] S. J. Brodsky and H. J. Lu, Phys. Rev. Lett. **64**, (1990) 1342.
- [49] M. Gyulassy and P. Levai, hep-ph/9807247.
- [50] A. Dumitru, D. H. Rischke, nucl-th/9806003.
- [51] X.-N. Wang, Phys. Rev. D **46**, R1900 (1992); Phys. Rev. D **47**, 2754 (1993).

